

Signs of Magnetic Accretion in the X-ray Pulsar Binary GX 301–2

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ABSTRACT

Observations of the cyclotron resonance scattering feature in the X-ray spectrum of GX 301–2 suggest that the surface field of the neutron star is $B_{\text{CRSF}} \sim 4 \times 10^{12}$ G. The same value has been derived in modelling the rapid spin-up episodes in terms of the Keplerian disk accretion scenario. However, the spin-down rate observed during the spin-down trends significantly exceeds the value expected in currently used spin-evolution scenarios. This indicates that either the surface field of the star exceeds $50 B_{\text{CRSF}}$, or a currently used accretion scenario is incomplete. We show that the above discrepancy can be avoided if the accreting material is magnetized. The magnetic pressure in the accretion flow increases more rapidly than its ram pressure and, under certain conditions, significantly affects the accretion picture. The spin-down torque applied to the neutron star in this case is larger than that evaluated within a non-magnetized accretion scenario. We find that the observed spin evolution of the pulsar can be explained in terms of the magnetically controlled accretion flow scenario provided the surface field of the neutron star is $\sim B_{\text{CRSF}}$.

Subject headings: X-rays: individual (GX 301-2/4U 1223-62/Wray 977) — stars: neutron — stars: magnetic fields — X-rays: binaries: close

1. Introduction

A strong magnetization of the neutron star in the High Mass X-ray Binary (HMXB) GX 301–2 has first been suspected by Lipunov (1982). Assuming the star to rotate at the equilibrium period given by Davidson & Ostriker (1973) he limited its dipole magnetic moment to $\gtrsim 10^{32}$ G cm³. This implies the surface field of the star to be in excess of 2×10^{14} G.

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This estimate has been challenged by observations of the cyclotron resonance scattering feature (CRSF) at $\sim 35 - 45$ keV (see La Barbera et al. 2005, and references therein). The field strength of the pulsar inferred from these observations, $B_{\text{CRSF}} \sim 4 \times 10^{12}$ G, is a factor of 50 smaller than the surface field evaluated by Lipunov (1982).

However, a question about the field strength of GX 301–2 has recently been raised again by Doroshenko et al. (2010) who reported the pulsar to spin down on a time scale of years at the average rate of $\dot{P}_s \simeq 4.25 \times 10^{-8} \text{ s s}^{-1}$ ($\dot{\nu}_{\text{sd}} = \dot{P}_0/P_s^2 \sim -10^{-13} \text{ Hz s}^{-1}$). Using the spin-down scenarios of Davidson & Ostriker (1973) and Illarionov & Kompaneets (1990) they put the same lower bound on the surface field of the neutron star as Lipunov (1982). To avoid the apparent discrepancy they have suggested that CRSF is generated at the top of a radiative-dominated accretion column, which is extended above the stellar surface up to an altitude of a few stellar radii.

In this paper we argue that this conclusion is premature. The spin behavior of the pulsar is very complicated. The long-term spin-down trends are superposed with phases of stable rotation and rapid spin-up episodes (see Figure 1, and Sect. 2). The value of torques applied to the star strongly depends on the geometry and physical conditions in the accreting material (Sect. 3). The pulsar behavior during the rapid spin-up episodes can be explained in terms of accretion from a Keplerian disk (Koh et al. 1997) provided the surface field of the star is $\sim B_{\text{CRSF}}$ (Sect. 4). The spin-down rate of the star during the spin-down trends strongly depends on the conditions in the material surrounding the pulsar magnetosphere (Sect. 5). We find that the neutron star can brake harder if the accreting material is magnetized. The accretion flow in this case is approaching the magnetosphere of the neutron star in a form of a dense magnetic slab confined by the magnetic field of the flow itself. The interaction between the slab and stellar magnetosphere leads to an effective angular momentum transfer from the star to the slab (Sect. 6). We show that the observed spin-down rate can be explained within this magnetically controlled accretion scenario provided the surface field of the neutron star is about B_{CRSF} (Sect. 7). Assumptions adopted in this accretion picture are briefly discussed in Sect. 8 and our conclusions are summarized in Sect. 9.

2. Basic parameters

GX 301–2 is a 685 s accreting X-ray pulsar in a 41.5 d eccentric orbit around the early B-type supergiant companion Wray 977 (Sato et al. 1986). The massive companion underfills its Roche lobe and loses material at the rate $\dot{M}_{\text{out}} \simeq 10^{-5} M_{\odot} \text{ yr}^{-1}$ in the form of a relatively slow ($v_w \sim 300 - 400 \text{ km s}^{-1}$) stellar wind (Kaper et al. 2006).

As the neutron star moves through the wind of density ρ_{∞} with a relative velocity $\mathbf{v}_{\text{rel}} = \mathbf{v}_{\text{ns}} + \mathbf{v}_w$ it captures material at a rate $\dot{\mathfrak{M}} \leq \dot{\mathfrak{M}}_c$, where

$$\dot{\mathfrak{M}}_c = \pi R_G^2 \rho_{\infty} v_{\text{rel}}. \quad (1)$$

Here R_{ns} is the radius, M_{ns} the mass and $v_{\text{ns}} \sim 250 \text{ km s}^{-1}$ is the orbital velocity of the neutron

star, and $R_G = 2GM_{\text{ns}}/v_{\text{rel}}^2$ is its Bondi radius.

The interaction between the accreting material and the stellar magnetic field leads to formation of a magnetosphere which, in the first approximation, prevents the accretion flow from reaching the stellar surface. The flow is decelerated by the stellar magnetic field at a distance r_m , which is defined by equating the pressure of the accretion flow with the magnetic pressure due to the dipole field of the neutron star and is referred to as the radius of the magnetosphere. The accretion flow enters the pulsar field at the magnetospheric boundary and flowing along the field lines reaches the surface of the neutron star at the magnetic pole regions. The rate at which the accreting material enters the pulsar field, $\dot{\mathfrak{M}}_{\text{in}}$, and, correspondingly, the mass accretion rate onto the stellar surface, $\dot{\mathfrak{M}}_{\text{a}}$, can be evaluated as

$$\dot{\mathfrak{M}}_{\text{in}} = \dot{\mathfrak{M}}_{\text{a}} \simeq 10^{17} R_6 m^{-1} \left(\frac{L_X}{2 \times 10^{37} \text{ erg s}^{-1}} \right) \text{ g s}^{-1}, \quad (2)$$

where $R_6 = R_{\text{ns}}/10^6 \text{ cm}$, $m = M_{\text{ns}}/1.4 M_{\odot}$ and L_X is the X-ray luminosity of the system, which is normalized to its average value $L_X \sim (1-3) \times 10^{37} \text{ erg s}^{-1}$ for a distance of 1.8–3 kpc, (Chichkov et al. 1995; Kaper et al. 2006).

The persistent character of the pulsar indicates that the magnetospheric radius of the neutron star is smaller than its corotation radius,

$$r_{\text{cor}} = \left(\frac{GM_{\text{ns}}}{\omega_s^2} \right)^{1/3} \simeq 1.3 \times 10^{10} m^{1/3} P_{685}^{2/3} \text{ cm}, \quad (3)$$

and hence, the centrifugal barrier at the magnetospheric boundary does not prevent the accretion flow from reaching the stellar surface (Shvartsman 1970). Here P_{685} is the spin period of the neutron star in units of 685 s, and $\omega_s = 2\pi/P_s$ is its angular velocity.

The pulsar shows a complicated picture of spin variation (see Figure 1), which can be roughly divided into three phases. The spin-down trends with the average rate $|\dot{\nu}_{\text{d0}}| \sim 10^{-13} \text{ Hz s}^{-1}$ and 1–3 years duration (Doroshenko et al. 2010) are superposed with the short-term (about 30 days) rapid spin-up episodes with $\dot{\nu}_{\text{u0}} \simeq 5 \times 10^{-12} \text{ Hz s}^{-1}$ (Koh et al. 1997), and the phase of almost stable rotation ($\dot{\nu} < 10^{-13} \text{ Hz s}^{-1}$), which lasts up to a few years (Bildsten et al. 1997). The pulsar transition between these phases occurs without any significant changes in the intensity and spectral parameters of the X-ray emission. In particular, X-ray luminosity of the system during rapid spin-up episodes is only a factor of 2–4 larger than that observed during the spin-down trends (Koh et al. 1997). This indicates that the mass-transfer between the system components operates almost stationary (i.e. $\dot{\mathfrak{M}}_{\text{c}} \sim \dot{\mathfrak{M}}_{\text{a}} = \dot{\mathfrak{M}}$) without any significant changes in the parameters of the accretion flow.

3. Accretion-driven spin evolution

The equation governing the spin evolution of an accreting neutron star reads

$$I\dot{\omega}_s = K_{\text{su}} - K_{\text{sd}}. \quad (4)$$

Here I is the moment of inertia of the neutron star, and K_{su} and K_{sd} are the spin-up and spin-down torques applied to the star from the accretion flow.

3.1. Spin-up torque

Since the material captured by the neutron star in a binary system possesses specific angular momentum, the mass accretion onto the neutron star is accompanied by the accretion of angular momentum at the rate $\dot{J} = \xi \Omega_{\text{orb}} R_G^2 \dot{\mathfrak{M}}_c$ (for discussion see, Illarionov & Sunyaev 1975; Illarionov & Kompaneets 1990, and references therein). Here $\Omega_{\text{orb}} = 2\pi/P_{\text{orb}}$ is the orbital angular velocity and ξ is a parameter accounting for dissipation of angular momentum in the accreting material (see e.g. Ruffert 1999, and references therein).

The geometry of the accretion flow in this case deviates from spherical symmetry. The deviation, however, will be significant only if $r_{\text{circ}} \gtrsim r_m$, where $r_{\text{circ}} = j^2/GM_{\text{ns}}\dot{\mathfrak{M}}_c^2$ is a circularization radius, at which the angular velocity of the accreting material, $\omega_{\text{en}} = \xi \Omega_{\text{orb}} (R_G/r)^2$, reaches the Keplerian angular velocity, $\omega_k = (r^3/2GM_{\text{ns}})^{1/2}$ (Bisnovatyi-Kogan 1991). If this condition is satisfied a formation of the Keplerian disk can be expected. Otherwise, the accretion should be treated in the Quasi-Spherical (QS) approximation, which implies $0 < \omega_{\text{en}}(r_m) \ll \omega_k(r_m)$.

The radius of the magnetosphere of a neutron star accreting from a non-magnetized wind (validity of this approximation is discussed in Sect. 6) is $r_m \sim r_a$, where (see, e.g. Arons 1993, and references therein)

$$r_a = \left(\frac{\mu^2}{\dot{\mathfrak{M}}(2GM_{\text{ns}})^{1/2}} \right)^{2/7}. \quad (5)$$

Here μ is the dipole magnetic moment of the neutron star. Solving inequality $r_{\text{circ}} \gtrsim r_a$ for v_{rel} one finds that the neutron star in a wind-fed HMXB would be surrounded by a persistent Keplerian disk if $v_{\text{rel}} \lesssim v_{\text{cr}}$, where (see e.g., Ikhsanov 2007),

$$v_{\text{cr}} \simeq 200 \xi_{0.2}^{1/4} \mu_{30}^{-1/14} m^{11/28} \dot{\mathfrak{M}}_{17}^{1/28} \left(\frac{P_{\text{orb}}}{41.5 \text{ d}} \right)^{-1/4} \text{ km s}^{-1}. \quad (6)$$

Here $\mu_{30} = \mu/10^{30} \text{ G cm}^3$, $\dot{\mathfrak{M}}_{17} = \dot{\mathfrak{M}}/10^{17} \text{ g s}^{-1}$ and $\xi_{0.2} = \xi/0.2$ is normalized to its average value derived in numerical modeling of quasi-spherical accretion in the non-magnetized flow approximation (Ruffert 1999, and references therein). The specific angular momentum of the material in the Keplerian disk is $j_k(r) = \omega_k r^2$. Therefore, the spin-up torque associated with the accretion of material from the inner radius of the disk is $K_{\text{su}}^{(\text{d})} = \dot{\mathfrak{M}} j_k(r_m) r_m^2 = (GM_{\text{ns}} r_m)^{1/2}$ (Pringle & Rees 1972).

Thus, the spin-up torque applied to an accreting neutron star in a wind-fed HMXB can be expressed as

$$K_{\text{su}} \simeq \begin{cases} \dot{\mathfrak{M}} (GM_{\text{ns}} r_{\text{m}})^{1/2} & \text{for } v_{\text{rel}} \leq v_{\text{cr}} \quad (\text{disk}) \\ \dot{\mathfrak{M}} \xi \Omega_{\text{orb}} R_{\text{G}}^2 & \text{for } v_{\text{rel}} > v_{\text{cr}} \quad (\text{QS}) \end{cases} \quad (7)$$

3.2. Spin-down torque

The spin-down torque applied to an accreting neutron star is associated with interaction between its magnetic field and material surrounding the magnetosphere. This interaction leads to a distortion of the magnetospheric field, turbulization of the material at the magnetospheric boundary, excitation of MHD waves and magnetic reconnection. The efficiency of these processes depends on the field geometry, conductivity of the accreting material and spectrum of the turbulent motions. A detailed modeling of these processes is, therefore, very complicated and is beyond the scope of this paper. Here we focus on a simplified task targeting at evaluation of upper limits to the spin-down torque under various assumptions about the geometry and parameters of the accretion flow.

The spin-down torque applied to a neutron star accreting material from a Keplerian disk has first been evaluated by Lynden-Bell & Pringle (1974). The stellar magnetosphere in this case contains both closed and open field lines, which are extended beyond the inner radius of the disk. The angular velocity of the material in those parts of the Keplerian disk, which are located inside the corotation radius, exceeds the angular velocity of the star itself. The interaction between these parts of the disk and stellar magnetosphere tends, therefore, to spin-up the star. The angular velocity of the material in the Keplerian disk is, however, decreasing with radius and becomes smaller than the angular velocity of the neutron star in the region $r > r_{\text{cor}}$. The interaction between the magnetospheric field lines and the disk at this distance tends to spin-down the star. As shown by Lynden-Bell & Pringle (1974), the spin-down torque associated with this interaction can be evaluated as $K_{\text{sd}}^{(\text{d})} = k_{\text{t}} \mu^2 / r_{\text{cor}}^3$, where $k_{\text{t}} < 1$ is the efficiency parameter.

The magnetosphere of a neutron star undergoing spherical accretion is closed. The stellar field is screened by the surface currents at the magnetospheric boundary and there is no field lines extended beyond the magnetospheric radius. The boundary is convex towards the accreting material and is closed to the cusp points located at the stellar magnetic axis (for discussion see Arons & Lea 1976; Elsner & Lamb 1977; Michel 1977). This indicates that the interaction between the stellar magnetosphere and the accretion flow occurs solely at the magnetospheric boundary. For the neutron star in this case to spin-down the angular velocity of the accreting material at the magnetospheric boundary, $\omega_{\text{en}}(r_{\text{m}})$, should be smaller than the angular velocity of the star itself, ω_{s} (Bisnovatyi-Kogan 1991). We assume here that this condition is satisfied and discuss the validity of this assumption in Sect 5.

The spin-down torque applied to a neutron star accreting from a quasi-spherical accretion flow,

in the first approximation, can be evaluated by finding the torque on a rotating sphere of the radius r_m in a viscous medium (for discussion see e.g., Lipunov 1992). The spin-down torque in this case can be expressed as

$$K_{sd} = 4\pi r_m^2 \nu_t \rho_0 v_\phi, \quad (8)$$

where ν_t is the viscosity coefficient, ρ_0 is the density of material at the magnetospheric boundary and $v_\phi = [\omega_s - \omega_{en}(r_m)] r_m$ is the azimuthal component of the relative linear velocity between the sphere and the surrounding material. For a simplicity we consider the case of $\omega_{en}(r_m) \ll \omega_s$ and hence, $v_\phi = \omega_s r_m$. This approximation is reasonable for evaluation of upper limits to the spin-down torque.

Assuming a turbulent nature of the viscosity, $\nu_t = k_t v_t \ell_t$, and taking into account that the density of the free-falling material at the magnetospheric boundary of the neutron star is $\rho_0 = \dot{\mathfrak{M}}_c / 4\pi r_m^2 v_{ff}(r_m)$, one finds

$$K_{sd} = k_t \omega_s r_m \ell_t(r_m) \dot{\mathfrak{M}}_c \frac{v_t(r_m)}{v_{ff}(r_m)}. \quad (9)$$

Here $v_t(r_m)$ and $\ell_t(r_m)$ are the velocity and the scale of turbulent motions, $v_{ff}(r_m) = (2GM_{ns}/r_m)^{1/2}$ is the free-fall velocity at the magnetospheric boundary and $k_t \leq 1$ is the efficiency parameter. If the turbulent motions in the material at the magnetospheric boundary are excited by its interaction with the rotating magnetosphere, the viscosity coefficient can be expressed as $\nu_t^{(0)} = k_t v_\phi r_m$. This implies that the scale and the velocity of turbulent motions are limited to $\ell_t \leq r_m$ and $v_t \leq v_\phi$. Putting these parameters to Eq. (9) and setting $r_m = r_a$ one finds

$$K_{sd}^{(0)} = k_t \frac{\mu^2 \omega_s^2}{GM_{ns}} = k_t \frac{\mu^2}{r_{cor}^3}. \quad (10)$$

This upper limit to the spin-down torque has previously been evaluated by Lipunov (1982) and Bisnovatyi-Kogan (1991) for the case of quasi-spherical accretion.

The neutron star under the same conditions is braking harder if $\nu_t > \nu_t^{(0)}$. This can be realized if $v_t > v_\phi$ (note, that the scale parameter, ℓ_t , has already been limited above to its maximum possible value). On the other hand, the velocity of turbulent motions is limited to the speed of sound in the accreting material since the supersonic turbulence it effectively suppresses by the Landau damping (for discussion see, e.g., Davies & Pringle 1981). Therefore, the maximum possible value of ν_t can be found by setting $v_t \sim v_{ff}$, i.e. assuming that the velocity of the turbulent motions in the accretion flow is close to the speed of sound in the material heated in the shock at the magnetospheric boundary up to the adiabatic temperature (for discussion see, e.g. Arons & Lea 1976; Elsner & Lamb 1977). Putting this to Eq. (9) yields

$$K_{sd}^{(t)} = k_t \dot{\mathfrak{M}}_c \omega_s r_m^2. \quad (11)$$

Finally, taking into account that $\dot{\mathfrak{M}} = \mu^2 / (2r_m^7 GM_{ns})^{1/2}$ (see Eq. 5), one can express the maximum possible spin-down torque as

$$K_{sd}^{(t)} = k_t \frac{\mu^2}{(r_m r_{cor})^{3/2}}. \quad (12)$$

The condition $v_t > v_\phi$ can be satisfied if the neutron star accretes material from either a hot turbulent atmosphere (as it occurs in the subsonic propeller scenario of Davies & Pringle 1981), or a magnetized flow with a high value of the magnetic viscosity. The subsonic propeller scenario can be used, however, for interpretation of a limited number of X-ray pulsars. The spin period of subsonic propellers is limited to $P_s < P_{br}$, where (Ikhsanov 2001)

$$P_{br} \simeq 450 \mu_{30}^{16/21} \mathfrak{M}_{15}^{-5/7} m^{-4/21} \text{ s} \quad (13)$$

is a so-called break period. If the spin period of a neutron star exceeds P_{br} the cooling of the atmosphere due to bremsstrahlung emission and turbulent motions becomes more effective than heating (due to propeller action by the neutron star). The X-ray source in this case would be observed as an X-ray burster (Lamb et al. 1977). Furthermore, the X-ray luminosity of subsonic propellers is limited to $L_X < L_{br}$, where

$$L_{br} \simeq 10^{34} \mu_{30}^{8/7} m^{-2/7} \mathfrak{M}_{15}^{-4/7} \left(\frac{P_s}{500 \text{ s}} \right)^{-1} \text{ erg s}^{-1}. \quad (14)$$

Otherwise, cooling of the material at the magnetospheric boundary due to inverse Compton scattering of the X-ray photons, emitted from the surface of the neutron star, on the hot atmospheric electrons dominates heating due to propeller action. This condition has been derived by solving inequality $t_c(r_m) > t_h$, where

$$t_c(r) = \frac{3\pi r^2 m_e c^2}{2 \sigma_T L_X} \quad (15)$$

is the Compton cooling time (Elsner & Lamb 1977) and $t_h \sim 1/\omega_s$ is the heating time due to propeller action (Davies & Pringle 1981). Here m_e is the electron mass and σ_T is the Thomson cross-section. Thus, the subsonic propellers can only appear as relatively low luminous, fast rotating pulsars, which is obviously not the case of GX 301–2.

A high viscosity can also be expected if the material surrounding the neutron star is magnetized. The magnetic viscosity coefficient in this case can be expressed as $\nu_m = k_t V_A \ell_m$, where $V_A = B_f/(4\pi\rho_0)^{1/2}$ is the Alfvén velocity and ℓ_m is the characteristic scale of the magnetic field, B_f , in the material at the magnetospheric boundary. A rapid braking of the neutron star in this case can be expected if $\ell_m \sim r_m$ and $V_A > V_\phi$. The Magnetically Controlled Accretion scenario in which this situation is realized will be discussed in Sect. 6.

Concluding this section one can express the spin-down torque applied to a neutron star from the accretion flow in the following form

$$K_{sd} = \begin{cases} k_t \frac{\mu^2}{r_{cor}^3} & \text{for } \nu_t \leq \nu_t^{(0)} \\ k_t \frac{\mu^2}{(r_m r_{cor})^{3/2}} & \text{for } \nu_t > \nu_t^{(0)} \end{cases} \quad (16)$$

where $k_t < 1$ is a parameter accounting for the conductivity, spectrum of turbulence and inhomogeneities of the accreting material.

4. Rapid spin-up episodes

The spin-up torque applied to the neutron star during the rapid spin-up episodes is $K_{\text{su}} = K_0 + K_{\text{sd}}$, where

$$K_0 = 2\pi I \dot{\nu}_{\text{su}} \simeq 3 \times 10^{34} I_{45} (\dot{\nu}_{\text{su}} / \dot{\nu}_{\text{u0}}) \text{ dyne cm}, \quad (17)$$

and $I_{45} = I / 10^{45} \text{ g cm}^2$. If the surface field of the neutron star does not considerably exceed B_{CRSF} the spin-down torque is significantly smaller than K_0 . In particular, the spin-down torque applied to the neutron star from the Keplerian disk under these conditions ($B_* = B_{\text{CRSF}}$) is $K_{\text{sd}}^{(\text{d})} = k_t \mu^2 / r_{\text{cor}}^3 \sim 2 \times 10^{30} k_t \text{ dyne cm}$ (see Eq. 16). The spin-up torque in this case can be evaluated as $K_{\text{su}} \gtrsim K_0$. Solving this inequality with spin-up torque expressed by Eq. (7) one finds that the spin behavior of the pulsar during the rapid spin-up episodes can be explained in terms of accretion from the Keplerian disk provided the dipole magnetic moment of the neutron star is $\mu \gtrsim \mu_0$, where

$$\mu_0 \simeq 2 \times 10^{30} \text{ G cm}^3 I_{45}^{7/2} \mathfrak{M}_{17}^{-3} m^{-3/2} \left(\frac{\dot{\nu}_{\text{su}}}{\dot{\nu}_{\text{u0}}} \right)^{7/2}. \quad (18)$$

This corresponds to the surface field $B_0 = 2\mu_0 / R_{\text{ns}}^3 \gtrsim 4 \times 10^{12} \text{ G}$, which is close to B_{CRSF} .

An attempt to explain the spin-up episodes in terms of the QS accretion scenario encounters difficulties. The spin-up torque in this case could be as high as K_0 only if the relative velocity of the neutron star were $v_{\text{rel}} \lesssim v_0$, where

$$v_0 \simeq 200 \xi_{0.2}^{1/4} m^{1/2} \mathfrak{M}_{17}^{1/4} \left(\frac{P_{\text{orb}}}{41.5 \text{ d}} \right)^{-1/4} \text{ km s}^{-1}. \quad (19)$$

However, a formation of the Keplerian disk under these conditions cannot be avoided (see Eq. 6). It therefore appears that the interpretation of the spin-up episodes in terms of the Keplerian disk accretion scenario is reliable and consistent with the surface field of the neutron star measured through observations of CRSF (for discussion see also Koh et al. 1997).

5. Spin-down trends

The spin-down torque applied to the neutron star during the spin-down trends is $|K_{\text{sd}}| = |K_1| + |K_{\text{su}}|$, where

$$|K_1| = 2\pi I \dot{\nu}_{\text{sd}} \simeq 6 \times 10^{32} I_{45} \left(\frac{\dot{\nu}_{\text{sd}}}{\dot{\nu}_{\text{d0}}} \right) \text{ dyne cm}. \quad (20)$$

An interpretation of the pulsar behavior during these trends cannot be made within the existing spin-down scenarios without additional assumptions. If the star accretes material from a Keplerian disk the spin-up torque is $|K_{\text{su}}^{(\text{d})}| \gg |K_1|$ (see above). The spin-down behavior of the pulsar in this case implies $|K_{\text{sd}}^{(\text{d})}| > |K_{\text{su}}^{(\text{d})}|$. To satisfy this condition one has to assume that the surface field of the star is in excess of $5 \times 10^{14} \text{ G}$, which is 100 times larger than B_{CRSF} .

The spin-up torque applied to the neutron star is significantly smaller if the star undergoes quasi-spherical accretion. Its value depends on the efficiency of the processes responsible for angular momentum dissipation in the accreting material which is accounted for by the parameter ξ (see Eq. 7). This parameter in the case under consideration is limited by the condition $\omega_{\text{en}}(r_{\text{m}}) < \omega_{\text{s}}$, which implies that the angular velocity of the accreting material at the magnetospheric boundary is smaller than the angular velocity of the star itself. Otherwise, the interaction between the accretion flow and stellar magnetosphere tends to spin-up the star (for discussion see, Bisnovatyi-Kogan 1991). Setting $r_{\text{m}} = r_{\text{a}}$ and solving the above inequality for the parameters of GX 301–2 yields

$$\begin{aligned} \xi < 0.03 \, m^{-12/7} L_{37}^{-4/7} R_6^{20/7} \left(\frac{P_{\text{orb}}}{41.5 \, \text{d}} \right) \times \\ \times \left(\frac{P_{\text{s}}}{685 \, \text{s}} \right)^{-1} \left(\frac{v_{\text{rel}}}{400 \, \text{km s}^{-1}} \right)^4 \left(\frac{B_{\text{*}}}{B_{\text{CRSF}}} \right)^{8/7} \end{aligned} \quad (21)$$

This indicates that $|K_{\text{su}}^{(0)}|$ during the spin-down trends does not exceed $|K_1|$ and hence, the condition for the spin-down phase can be expressed as $|K_{\text{sd}}| \gtrsim |K_1|$.

Putting $K_{\text{sd}} \sim K_{\text{sd}}^{(0)}$ one finds that the above condition can be satisfied only if the surface field of the neutron star is $\gtrsim 10^{14} \, k_{\text{t}}^{-1/2} \, \text{G}$ (for discussion see also Lipunov 1982; Doroshenko et al. 2010). To satisfy the condition $K_{\text{sd}}^{(\text{t})} \gtrsim K_1$ one has to assume that either the surface field of the neutron star is $\gtrsim 3 \, k_{\text{t}}^{-7/8} B_{\text{CRSF}}$, or the radius of the stellar magnetosphere is smaller than its canonical value, r_{a} , at least by a factor of $2 \, k_{\text{t}}^{-2/3}$.

Thus, both the Keplerian disk and quasi-spherical accretion scenarios encounters difficulties explaining the spin behavior of GX 301–2 during the spin-down trends. This may indicate that either the surface field of the star is indeed in excess of B_{CRSF} , or the neutron star in this binary system accretes in a different way. The first possibility has been already discussed by Doroshenko et al. (2010). Here we focus on the analysis of an alternative accretion scenario. We show that the peculiar spin behavior of GX 301–2 can be explained in terms of the magnetic accretion scenario (Shvartsman 1971; Bisnovatyi-Kogan & Ruzmaikin 1974, 1976) provided the material captured by the neutron star from the wind of its normal companion is magnetized.

6. Magnetic accretion

Let us consider a situation in which the magnetic energy density in the material captured by the neutron star, $\mathcal{E}_{\text{m}}(R_{\text{G}}) = B_{\text{f0}}^2/8\pi$, is comparable to its thermal energy density, $\mathcal{E}_{\text{th}}(R_{\text{G}}) = \rho_{\infty} c_{\text{s0}}^2$, i.e. $\beta \equiv \mathcal{E}_{\text{th}}/\mathcal{E}_{\text{m}} \sim 1$. Here c_{s0} is the speed of sound and B_{f0} is the strength of the magnetic field in the captured material at R_{G} . The magnetic field in the free-falling material is dominated by its radial component, B_{r} , (the transverse scales in the free-falling flow contract as r^{-2} , while the radial scales expand as $r^{1/2}$, Zeldovich & Shakura 1969), which under the magnetic flux conservation condition increases as $B_{\text{r}}(r) \sim B_{\text{f0}} (R_{\text{G}}/r)^2$ (Bisnovatyi-Kogan & Fridman 1970). The magnetic

pressure in the free-falling material,

$$\mathcal{E}_m(r) = \mathcal{E}_m(R_G) \left(\frac{R_G}{r} \right)^4, \quad (22)$$

increases, therefore, more rapidly than its ram pressure,

$$\mathcal{E}_{\text{ram}}(r) = \mathcal{E}_{\text{ram}}(R_G) \left(\frac{R_G}{r} \right)^{5/2}, \quad (23)$$

and hence, the gravitational energy of the star is converted predominantly into the magnetic pressure of the free-falling material, $\mathcal{E}_m/\mathcal{E}_{\text{ram}} \propto r^{-3/2}$. Here $\mathcal{E}_{\text{ram}}(R_G) = \rho_\infty v_{\text{rel}}^2$ is the ram pressure of the captured material at the Bondi radius.

The distance R_{sh} (hereafter Shvartsman radius) at which the magnetic pressure in the accretion flow reaches its ram pressure can be derived by solving equation $\mathcal{E}_m(R_{\text{sh}}) = \mathcal{E}_{\text{ram}}(R_{\text{sh}})$. This yields (Shvartsman 1971),

$$R_{\text{sh}} = \beta^{-2/3} \left(\frac{c_s}{v_{\text{rel}}} \right)^{4/3} R_G = \beta^{-2/3} \frac{2GM_{\text{ns}} c_s^{4/3}}{v_{\text{rel}}^{10/3}}. \quad (24)$$

If the magnetic flux in the accreting material is conserved, the accretion ends at the Shvartsman radius. Further accretion in this case is impossible. Otherwise, the magnetic energy in the flow would exceed its gravitational energy, which contradicts the energy conservation law (for discussion see Shvartsman 1971). Therefore, the accretion flow can approach the star to a closer distance only if dissipation of the magnetic field in the flow occurs. If the field dissipation is governed by magnetic reconnection the characteristic time of the accretion process inside R_{sh} is limited to $t \geq t_{\text{rec}}$, where

$$t_{\text{rec}} = \frac{r}{\eta_m v_A} = \eta_m^{-1} t_{\text{ff}} \left(\frac{v_{\text{ff}}}{v_A} \right). \quad (25)$$

The value of the efficiency parameter η_m depends on physical conditions and field configuration in the region of reconnection, and in the general case is limited to $0 < \eta_m \leq 0.1$ (see, e.g., Parker 1971; Kadomtsev 1987; Noglik et al. 2005; Somov 2006, and references therein). The Alfvén velocity in the accreting material increases as it approaches the neutron star from the initial value $V_{A0} \sim \beta^{-1/2} c_s$ to $V_A(R_{\text{sh}}) \sim V_{\text{ff}}$. Hence, the time of the magnetic reconnection in the accretion flow remains significantly larger than the dynamical (free-fall) time, $t_{\text{ff}} = r/V_{\text{ff}} = (r^3/2GM_{\text{ns}})^{1/2}$, at any stage of the accretion process. This proves reliability of the assumption about conservation of the magnetic flux in the free-falling material. But, on the other hand, it also suggests that the accretion flow will be decelerated at the Shvartsman radius by its own magnetic field and the accretion process in the region $r < R_{\text{sh}}$ operates in the diffusion approximation.

Rapid amplification of the magnetic field in the spherical flow and deceleration of the flow by its own magnetic field at the Shvartsman radius have been confirmed in analytical studies (Bisnovaty-Kogan & Ruzmaikin 1974, 1976) and numerical calculations of magnetized spherical accretion onto a black hole (Igumenshchev et al. 2003; Igumenshchev 2006). These calculations

have shown that the magnetized flow is shock-heated at the Shvartsman radius up to the adiabatic temperature. The structure of the accretion flow inside the Shvartsman radius depends on the efficiency of cooling of the accreting material. If cooling is inefficient the accretion process switches into the convective-dominated stage in which some material is leaving the system in a form of jets (Igumenshchev et al. 2003). Otherwise, the material tends to flow along the lines of the large scale field of the accreting matter itself and accumulates in a dense non-Keplerian slab (see Fig. 1 in Bisnovaty-Kogan & Ruzmaikin 1976). The material in the slab is confined by its own magnetic field and its radial motion continues as the field is annihilating. This indicates that the accretion process in the slab occurs on the reconnection timescale expressed by Eq. (25).

6.1. Magnetic accretion in X-ray pulsars

The Shvartsman radius exceeds the canonical magnetospheric radius of the neutron star evaluated in the quasi-spherical accretion scenario, r_a , if the relative velocity of the neutron star through the wind of its massive companion satisfies the condition $v_{\text{rel}} \leq v_{\text{mca}}$, where

$$v_{\text{mca}} = \beta^{-1/5} (2GM_{\text{ns}})^{12/35} \mu^{-6/35} \mathfrak{M}^{3/35} c_s^{2/5}. \quad (26)$$

For typical parameters of HMXBs this velocity is

$$v_{\text{mca}} \simeq 680 \beta^{-1/5} m^{12/35} \mu_{30}^{-6/35} \mathfrak{M}_{17}^{3/35} c_6^{2/5} \text{ km s}^{-1},$$

which under the conditions of interest substantially exceeds the critical velocity at which the formation of the Keplerian disk in the system can be expected, v_{cr} (see Eq. 6). Here $c_6 = c_{s0}/10^6 \text{ cm s}^{-1}$. This finding allows us to distinguish a subclass of HMXBs in which the accretion occurs in a spherically symmetrical fashion and the accreting material is strongly affected by the magnetic field of the flow itself. This subclass is defined by the condition $v_{\text{cr}} < v_{\text{rel}} < v_{\text{mca}}$.

The structure of the accretion flow inside the Shvartsman radius depends on the efficiency of cooling processes in the accreting material. As shown by Arons & Lea (1976) and Elsner & Lamb (1977), the cooling of the material accreting onto a neutron star is dominated by the inverse Compton scattering of X-ray photos emitted from the stellar surface on the hot electrons of the accretion flow. This mechanism in the considered case will be effective if the Compton cooling time of the accretion flow at the Shvartsman radius, $t_c(R_{\text{sh}})$, is smaller than the characteristic time of the accretion process, which in the case of the magnetic accretion scenario is the reconnection time, t_{rec} . Combining Eqs. (15) and (25), one can express inequality $t_c(R_{\text{sh}}) \leq t_{\text{rec}}$ as $L_X \gtrsim L_{\text{cr}}$, where

$$L_{\text{cr}} \simeq 3 \times 10^{33} \mu_{30}^{1/4} m^{1/2} R_6^{-1/8} \left(\frac{\eta_{\text{m}}}{0.001} \right) \left(\frac{R_{\text{sh}}}{r_a} \right)^{1/2} \text{ erg s}^{-1}. \quad (27)$$

This indicates that the cooling of magnetized flow can be effective even in faint X-ray pulsars in which the Shvartsman radius does not significantly exceed the Alfvén radius of the neutron star. The structure of the accretion flow in those systems in which the above conditions are satisfied can be treated in terms of the magnetic slab.

6.2. Magnetospheric radius

The material in the slab is approaching the neutron star up to a distance at which its gas pressure becomes equal to the magnetic pressure due to the stellar dipole magnetic field. The gas density in the slab at this distance (the magnetospheric radius) is, therefore,

$$\rho_{\text{sl}} = \frac{\mu^2 m_{\text{p}}}{2\pi k_{\text{B}} T_0 r_{\text{m}}^6}, \quad (28)$$

where T_0 is the temperature of the material at the inner radius of the slab.

The value of the magnetospheric radius in the general case depends on the mode by which the material enters the stellar field at the magnetospheric boundary. As shown by Elsner & Lamb (1984), the diffusion rate of the accreting material into the stellar field can be evaluated as

$$\begin{aligned} \dot{\mathcal{M}}_{\text{in}}(r_{\text{m}}) &= 4\pi r_{\text{m}} \delta_{\text{m}} \rho_0 V_{\text{ff}}(r_{\text{m}}) \\ &= 4\pi r_{\text{m}}^{5/4} D_{\text{eff}}^{1/2} \rho_0 (2GM_{\text{ns}})^{1/4}, \end{aligned} \quad (29)$$

where $\delta_{\text{m}} = (D_{\text{eff}} \tau_{\text{d}})^{1/2}$ is the thickness of the diffusion layer at the magnetospheric boundary (magnetopause) and D_{eff} is the effective diffusion coefficient. The diffusion time is determined by the time on which the material being penetrated into the field leaves the magnetopause by free-falling along the magnetospheric field lines towards the stellar surface, $\tau_{\text{d}} \sim t_{\text{ff}}(r_{\text{m}})$.

The estimate (29) has good theoretical and observational grounds. It has been justified by studies of the Earth’s magnetosphere which show that the rate of solar wind penetration into the magnetic field of the Earth can be evaluated taking $D_{\text{eff}} \sim D_{\text{B}}$, where

$$D_{\text{B}} = \alpha_{\text{B}} \frac{ck_{\text{B}} T_0 r_{\text{m}}^3}{2e\mu} \quad (30)$$

is the Bohm diffusion coefficient, e is the electron charge and α_{B} is the efficiency parameter, which ranges in $0.1 - 0.25$ (Gosling et al. 1991). The diffusion process in this case is governed by magnetic reconnection and drift-dissipative instabilities (Paschmann 2008, and references therein). Note also that the same diffusion rate has previously been measured in plasma experiments with TOKAMAKs (Kadomtsev & Shafranov 1983) and evaluated from observations of solar flares (Priest 1982).

A higher rate of plasma penetration into the magnetosphere could be expected if the magnetospheric boundary were interchange unstable. The Rayleigh-Taylor instability of the boundary can occur in bright long-period X-ray pulsars. This instability would be, however, suppressed by the magnetic pressure gradient if the X-ray luminosity of the neutron star undergoing spherical accretion is $L_{\text{X}} < 3 \times 10^{36} \text{ erg s}^{-1}$ (Arons & Lea 1976; Elsner & Lamb 1977), and by the magnetic field shear in the magnetopause if the spin period of the neutron star is smaller than a few hundred seconds (see e.g. Ikhsanov & Pustilnik 1996, and references therein). The Kelvin-Helmholtz instability can be effective in short-period X-ray pulsars in which the magnetospheric radius of the neutron star is close to its corotation radius (Burnard, Lea & Arons 1983). But this instability might not be

effective if the relative velocity between the magnetosphere and the accretion flow exceeds the speed of sound or/and if the material at the boundary is magnetized (for discussion see Anzer & Börner 1980, 1983; Malagoli, Bodo & Rosner 1996). This indicates that the interchange instabilities of the boundary can be at work only in a limited number of pulsars, while Bohm diffusion takes place in any of the considered objects. Having this in mind, we focus our consideration on the case of Bohm diffusion assuming that the interchange instabilities of the boundary are suppressed.

The stationary accretion picture implies that the rate of plasma diffusion into the pulsar field at the magnetospheric boundary is equal to the mass capture rate by the neutron star from its environment at the Bondi radius and to the mass accretion rate onto the surface of the neutron star. Let us assume that this condition is satisfied at the distance r_{mca} , i.e. $\dot{\mathfrak{M}}_{\text{in}}(r_{\text{mca}}) = L_{\text{X}} R_{\text{ns}} / GM_{\text{ns}}$. Combining Eqs. (28), (29) and (30), and solving the above equation for r_{mca} , one gets

$$r_{\text{mca}} \simeq 8 \times 10^7 \alpha_{0.1}^{2/13} \mu_{30}^{6/13} T_6^{-2/13} m^{5/13} R_6^{-4/13} L_{37}^{-4/13} \text{ cm} \quad (31)$$

Here $\alpha_{0.1} = \alpha/0.1$, $T_6 = T_0/10^6 \text{ K}$ is the plasma temperature at the magnetospheric boundary, which is normalized according to Masetti et al. (2006), and L_{37} is the X-ray luminosity of the pulsar in units of $10^{37} \text{ erg s}^{-1}$.

The value of r_{mca} under the same conditions is smaller than the value of r_{a} , which represents the canonical radius of the magnetosphere of the neutron star undergoing spherical accretion. This can be associated with accumulation of material at the inner radius of the slab which occurs if $\dot{\mathfrak{M}}_{\text{in}} < \dot{\mathfrak{M}}_{\text{c}}$. As the material accumulates the gas pressure increases and the slab approaches the neutron star to a closer distance. Since $\dot{\mathfrak{M}}_{\text{in}} \propto r_{\text{m}}^{-13/4}$ the diffusion rate of material into the magnetosphere also increases and reaches $\dot{\mathfrak{M}}_{\text{c}}$ as the inner radius of the slab decreases to r_{mca} .

The spin-down torque applied to the neutron star from the magnetic slab can be evaluated as $K_{\text{sd}}^{(\text{sl})} = k_{\text{t}} \dot{\mathfrak{M}} \omega_{\text{s}} r_{\text{mca}}^2$. Taking into account that $\dot{\mathfrak{M}} = \mu^2 / (2GM_{\text{ns}} r_{\text{mca}}^7)^{1/2}$ one finds

$$K_{\text{sd}}^{(\text{sl})} = \frac{k_{\text{t}} \mu^2 \omega_{\text{s}}}{r_{\text{mca}}^{3/2} (2GM_{\text{ns}})^{1/2}} = \frac{k_{\text{t}} \mu^2}{(r_{\text{mca}} r_{\text{cor}})^{3/2}} \quad (32)$$

Thus, the spin-down torque applied to the neutron star accreting material from the magnetic slab is a factor of $\sim (r_{\text{a}}/r_{\text{mca}})^{3/2}$ higher than the maximum possible spin-down torque applied to the star accreting material at the same rate from a quasi-spherical flow.

7. Signs of magnetic accretion in GX 301–2

The maximum possible spin-down rate of the neutron star in GX 301–2 within the magnetically controlled accretion scenario can be evaluated as $\dot{\nu}_{\text{sd}}^{(\text{mca})} = K_{\text{sd}}^{(\text{sl})} / 2\pi I$, that for the pulsar parameters is

$$\dot{\nu}_{\text{sd}}^{(\text{mca})} \simeq 7 \times 10^{-13} \text{ Hz s}^{-1} k_{\text{t}} \alpha_{0.1}^{-3/13} m^{-14/13} \quad (33)$$

$$\times I_{45}^{-1} T_6^{3/13} L_{37}^{6/13} R_6^{57/13} P_{685}^{-1} \left(\frac{B_*}{B_{\text{CRSF}}} \right)^{17/13}$$

Hence, the spin-down rate of the neutron star observed during the spin-down trends can be explained in terms of the magnetically controlled accretion scenario provided the surface field of the star is close to B_{CRSF} and the efficiency coefficient is $k_t \geq 0.14$. The spin-down power of the neutron star within this scenario is spent in the energy of electric currents and turbulent motions excited at the magnetospheric boundary and, possibly, to spin-up the material at the inner radius of the magnetic slab. In this case the angular momentum lost by the star can partly be accumulated in the material surrounding its magnetosphere. The spin-down torque in this case decreases as the angular velocity of the material at the inner radius of the slab approaches the angular velocity of the star itself. The angular momentum accumulated in the slab can later be transferred back to the star due to accretion process. The star in this case is expected to switch from the spin-up to the spin-down phase without any significant change of the accretion flow geometry.

A detailed study of the angular momentum exchange between the star and the slab is beyond the scope of this paper. Here we would like to note only that the angular velocity to which the material in the slab could be spun-up within this scenario is limited to the angular velocity of the star itself, ω_s . The angular momentum transfer from the star to the slab can bring its inner part to a solid rotation. The radial size of this part is limited, however, to the corotation radius at which $\omega_s = \omega_k(r_{\text{cor}})$. This opens a possibility to consider a situation in which the spin-up of the material at the inner radius of the slab results in temporary transition of the accretion picture from the non-Keplerian slab to Keplerian disk of the radius r_{cor} . The star in this case will experience a rapid spin-up on a timescale of the viscous time at the corotation radius. It is interesting that the size of the Keplerian disk evaluated by Koh et al. (1997) from modelling the rapid spin-up events is very close to the corotation radius of the pulsar.

8. Discussion

One of the basic assumptions adopted in the magnetically controlled accretion scenario is that the magnetic field energy density in the material captured by the neutron star at the Bondi radius is comparable to its thermal energy density. The latter can be evaluated by taking into account that $\rho_\infty = \dot{\mathfrak{M}}/\pi R_G^2 v_{\text{rel}}$. This yields

$$\mathcal{E}_{\text{th}} \simeq 0.02 \text{ erg cm}^{-3} m^{-2} c_6^2 \dot{\mathfrak{M}}_{17} \left(\frac{v_{\text{rel}}}{500 \text{ km s}^{-1}} \right)^3. \quad (34)$$

The magnetic energy density in the captured wind depends on the field strength of the massive companion of the neutron star. The dipole approximation ($B_{\text{mc}} \propto a^{-3}$) to the magnetic field of the massive stars remains valid up to a distance a_k , at which the dynamical pressure of the wind ejecting by this star reaches the magnetic tension of the stellar dipole field (here a is the distance from the massive companion). The magnetic field in the material propagating beyond a_k decreases

as $B \propto a^{-2}$ (Walder, Folini & Meynet 2011). Therefore, the magnetic energy density in the wind at the binary separation $a_0 > a_k$ is $\mathcal{E}_m(a_0) = \mathcal{E}_{m0} = \mu_{ms}^2 / (2\pi a_k^2 a_0^4)$, which implies

$$\mathcal{E}_{m0} \simeq 0.33 \text{ erg cm}^{-3} a_{13}^{-4} \left(\frac{\mu_{ms}}{10^{39} \text{ G cm}^3} \right)^2 \left(\frac{a_k}{100 R_\odot} \right)^{-2} \quad (35)$$

Here μ_{ms} is the dipole magnetic moment of the massive star and a_{13} is the binary separation in units of 10^{13} cm. It, therefore, appears, that the magnetic accretion scenario can be expected in GX 301–2 if the surface field of the massive component is a few hundred Gauss.

Recent spectropolarimetric observations (see, e.g. Hubrig et al. 2006; Oksala et al. 2010; Martins et al. 2010, and references therein) have shown a relatively strong magnetization of O/B-type stars to be not unusual. The strength of the large-scale field at the surface of these objects has been measured in the range $\sim 500 - 5000$ G, and in some cases beyond 10 kG. In this light, the assumption about relatively strong magnetization of the massive star in GX 301–2 seems to be rather reliable. It can be tested by spectropolarimetric observations of this system, which, therefore, appears to be of great importance for justification of the magnetically controlled accretion scenario used in our paper.

9. Conclusion

We have shown that the rapid spin-down of the neutron star observed during the spin-down trends of GX 301–2 can be explained in terms of magnetically controlled accretion scenario provided the surface field of the star is close to the value, B_{CRSF} , derived from observations of the cyclotron resonance scattering feature in the X-ray spectrum of this pulsar. The same strength of the magnetic field has been evaluated from modeling of the rapid spin-up events in terms of the Keplerian disk accretion. We conclude that the problems in modeling of the spin evolution of GX 301–2 encountered by Doroshenko et al. (2010) indicate an oversimplification in the currently used accretion scenarios rather than an extremely high magnetization of the neutron star. We have shown that these problems can be avoided by incorporation of the magnetic field of the accretion flow into the quasi-spherical accretion scenario provided the relative velocity of the neutron star satisfies the condition $v_{\text{cr}} < v_{\text{rel}} < v_{\text{mca}}$. The accretion flow in this case is decelerated by its own magnetic field at the Shvartsman radius, which exceeds the magnetospheric radius of the neutron star, and accumulates in the non-Keplerian magnetic slab confined by the magnetic field of the flow itself. The plasma approaches the star on the timescale of field dissipation in the magnetic slab, which significantly exceeds the dynamical (free-fall) timescale. The process of plasma penetration into the stellar magnetosphere is governed by the magnetic reconnection and drift-dissipative instabilities and occurs at a rate of the Bohm diffusion. The magnetospheric radius within this scenario, r_{mca} , is smaller than the canonical value, r_a . The spin-down torque applied to the neutron star from the slab is $K_{\text{sd}} = k_t \mu^2 / (r_{\text{mca}} r_{\text{cor}})^{3/2}$, where $0.1 < k_t < 1$. An exchange of the angular momentum between the star and the slab can lead to variations of the spin period of the pulsar.

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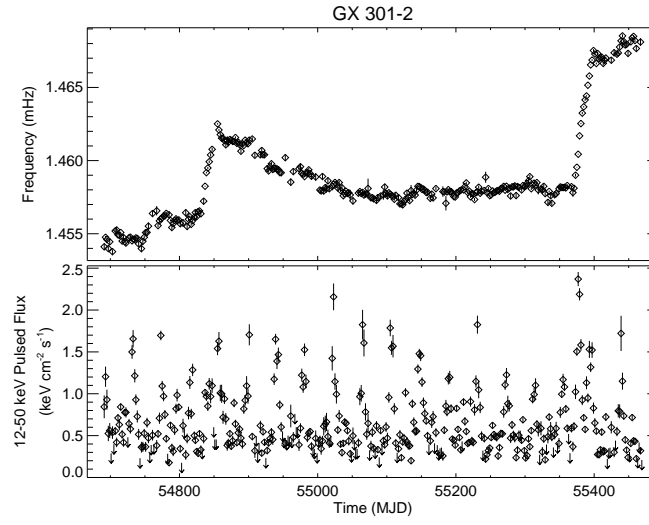


Fig. 1.— Pulse frequency and 12–50 keV pulsed flux observed with the Fermi Gamma-ray Burst Monitor detectors. GMB Pulsar Project <http://gammaray.nsstc.nasa.gov/gbm/science/pulsars/>